

Dark matter and Higgs phenomenology predicted by left-right twin Higgs model in light of CDMS II results

Lei Wang¹, Jin Min Yang²

¹ *Department of Physics, Yantai University, Yantai 264005, PR China*

² *Key Laboratory of Frontiers in Theoretical Physics,
Institute of Theoretical Physics, Academia Sinica, Beijing 100190, PR China*

Abstract

The left-right twin Higgs model predicts a light stable scalar \hat{S} , which is a candidate for WIMP dark matter. We study its scattering on nucleon and find that the cross section is below the CDMS II upper bound but can reach the SuperCDMS sensitivity. Then we study the Higgs phenomenology by paying special attention to the decay $h \rightarrow \hat{S}\hat{S}$ which is strongly correlated with the dark matter scattering on nucleon. We find that such an invisible decay can be sizable, which can severely suppress the conventional decay modes like $h \rightarrow VV$ ($V = W, Z$) and $h \rightarrow b\bar{b}$. On the other hand, compared to the SM prediction, the rates of Higgs boson productions at the LHC via gluon-gluon fusion, weak boson fusion or in association with top quark pairs are all reduced significantly, e.g., the gluon-gluon fusion channel can be suppressed by about 30%.

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I. INTRODUCTION

The twin Higgs mechanism [1, 2] is proposed as an interesting solution to the hierarchy problem. The SM Higgs emerges as a pseudo-Goldstone boson once a global symmetry is spontaneously broken, which is similar to what happens in the little Higgs models [3]. An additional discrete symmetry is imposed, which ensures the absence of one-loop quadratic divergence of Higgs mass. The resulting Higgs boson mass is naturally around the electroweak scale when the cut-off scale of the theory is around 5-10 TeV. The twin Higgs mechanism can be implemented in left-right models with the additional discrete symmetry being identified as the left-right symmetry [2]. In the left-right twin Higgs (LRTH) model, several physical Higgs bosons remain after the spontaneous symmetry breaking. Another additional discrete symmetry is introduced in the model under which the $SU(2)_L$ doublet \hat{h} is odd while all the other fields are even. The lightest particle \hat{S} in its neutral components is stable and thus can be a candidate for weakly interacting massive particle (WIMP) dark matter. The phenomenology of LRTH model has been studied by some authors [4].

The density of cold dark matter in the universe has been determined precisely by WMAP [5]:

$$\Omega_{CDM}h^2 = 0.105^{+0.021}_{-0.030}. \quad (1)$$

The thermal production of WIMPs can naturally explain such a relic density. As a direct detection of WIMPs, the CDMS attempts to observe the recoil energy transferred to a target nucleus in an elastic collision with a WIMP. Very recently the CDMS collaboration has completed their analysis of the final data runs of the CDMS II experiment and reported two candidate events [6]. Although these events cannot be interpreted as significant evidence for WIMP interacting with nucleons, the CDMS gives the most stringent upper limit on the WIMP-nucleon spin-independent cross section. For example, the cross section is constrained to be smaller than $3.8 \times 10^{-44} \text{ cm}^2$ for a WIMP of 70 GeV at 90% confidence level [6]. The implications of the new results from the CDMS II experiment have been discussed in many models [7].

In this work we focus on the left-right twin Higgs model. We first examine the scattering of the dark matter candidate \hat{S} with nucleon and compare the rate with the CDMS II results. Then we study the Higgs phenomenology, paying special attention to the decay $h \rightarrow \hat{S}\hat{S}$ which is strongly correlated with the dark matter scattering on nucleon. We will figure

out the size of such an invisible decay rate and how severely to suppress the conventional decay modes like $h \rightarrow VV$ ($V = W, Z$) and $h \rightarrow b\bar{b}$. We also study the suppression for the rates of Higgs boson productions at the LHC via gluon-gluon fusion, weak boson fusion or in association with a pair of top quarks. Since the LHC will be able to discover the Higgs boson in the full mass range [8], our study will help to probe the left-right twin Higgs model.

This work is organized as follows. In Sec. II, we briefly review the left-right twin Higgs model. In Sec. III, we examine the scattering of the dark matter candidate \hat{S} with nucleon and compare the rate with the CDMS II results. Also, the correlation of Higgs decays with the dark matter scattering on nucleon is studied. In Sec. IV, we calculate the main productions of the Higgs boson at the LHC. Finally, we give our conclusion in Sec. V.

II. LEFT-RIGHT TWIN HIGGS MODEL

A. Mass terms of gauge bosons

In LRTH model [2, 9], the global symmetry is $U(4) \times U(4)$ with a gauged $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ subgroup. The twin symmetry is identified as a left-right symmetry which interchanges L and R , implying that that gauge couplings of $SU(2)_L$ and $SU(2)_R$ are identical ($g_{2L} = g_{2R} = g_2$).

A pair of Higgs fields, H and \hat{H} , are introduced and each transforms as $(\mathbf{4}, \mathbf{1})$ and $(\mathbf{1}, \mathbf{4})$ respectively under the global symmetry. They can be written as

$$H = \begin{pmatrix} H_L \\ H_R \end{pmatrix}, \quad \hat{H} = \begin{pmatrix} \hat{H}_L \\ \hat{H}_R \end{pmatrix}, \quad (2)$$

where $H_{L,R}$ and $\hat{H}_{L,R}$ are two component objects which are charged under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ as

$$H_L \text{ and } \hat{H}_L : (\mathbf{2}, \mathbf{1}, 1), \quad H_R \text{ and } \hat{H}_R : (\mathbf{1}, \mathbf{2}, 1). \quad (3)$$

Each Higgs acquires a non-zero VEV as

$$\langle H \rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ f \end{pmatrix}, \quad \langle \hat{H} \rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ \hat{f} \end{pmatrix}, \quad (4)$$

which breaks one of the $U(4)$ to $U(3)$ and yields seven Nambu-Goldstone bosons. The scalar fields can be parameterized as

$$H = f e^{i\frac{\pi}{f}} \begin{pmatrix} 0 \\ 0 \\ 0 \\ 1 \end{pmatrix}, \quad \text{with } \pi = \begin{pmatrix} -N/2 & 0 & 0 & h_1 \\ 0 & -N/2 & 0 & h_2 \\ 0 & 0 & -N/2 & C \\ h_1^* & h_2^* & C^* & 3N/2 \end{pmatrix}, \quad (5)$$

with π being the corresponding Goldstone fields. N is a neutral real pseudoscalar, C and C^* are a pair of charged complex scalar fields, and $(h_1, h_2)^T$ is the SM $SU(2)_L$ Higgs doublet. \hat{H} can be parameterized in the same way by its own Goldstone fields $\hat{\pi}$, which contains \hat{N} , \hat{C} and $\hat{h} = (\hat{h}_1^+, \hat{h}_2^0)^T$.

The generators of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ are given respectively as

$$\begin{pmatrix} \frac{1}{2}\sigma_i & 0 \\ 0 & 0 \end{pmatrix}, \quad \begin{pmatrix} 0 & 0 \\ 0 & \frac{1}{2}\sigma_i \end{pmatrix}, \quad \frac{1}{2} \begin{pmatrix} 1_2 & 0 \\ 0 & 1_2 \end{pmatrix}, \quad (6)$$

and the corresponding gauge fields are

$$W_2 = \frac{1}{2} \begin{pmatrix} W_L^0 & \sqrt{2}W_L^+ & 0 & 0 \\ \sqrt{2}W_L^- & -W_L^0 & 0 & 0 \\ 0 & 0 & W_R^0 & \sqrt{2}W_R^+ \\ 0 & 0 & \sqrt{2}W_R^- & -W_R^0 \end{pmatrix}, \quad W_{B-L} = \frac{W_1}{2} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \quad (7)$$

where the Lorentz indices are suppressed. The covariant derivative is

$$D^\mu = \partial^\mu - ig_2 W_2^\mu - ig_1 n_{B-L} W_{B-L}^\mu, \quad (8)$$

where g_1 and g_2 are the gauge couplings for $U(1)_{B-L}$ and $SU(2)_{L,R}$, and n_{B-L} is the charge of the field under $U(1)_{B-L}$.

The covariant kinetic terms of Higgs fields can be written down as [2, 9]

$$\mathcal{L}_H = (D_\mu H)^\dagger D^\mu H + (D_\mu \hat{H})^\dagger D^\mu \hat{H}, \quad (9)$$

with $n_{B-L} = 1$. The above Lagrangian contains the following neutral Higgs boson interactions:

$$\begin{aligned} \mathcal{L}_H \supset & \frac{1}{2} g_2^2 f^2 s_1^2 W_L^- W_L^+ + \frac{1}{2} g_2^2 (\hat{f}^2 + f^2 c_1^2) W_R^- W_R^+ + \frac{1}{4} g_1^2 (f^2 + \hat{f}^2) W_1 W_1 \\ & - \frac{1}{4} g_1 g_2 f^2 (1 - c_2) W_1 W_L^0 + \frac{1}{8} g_2^2 f^2 (1 - c_2) W_L^0 W_L^0 - \frac{1}{4} g_1 g_2 (f^2 + f^2 c_2 + 2\hat{f}^2) W_1 W_R^0 \\ & + \frac{1}{8} g_2^2 (f^2 + f^2 c_2 + 2\hat{f}^2) W_R^0 W_R^0, \end{aligned} \quad (10)$$

where

$$\begin{aligned} c_1 &= \cos \frac{h+v}{\sqrt{2}f}, & s_1 &= \sqrt{1-c_1^2}, \\ c_2 &= \cos \frac{\sqrt{2}(h+v)}{f}, & s_2 &= \sqrt{1-c_2^2}. \end{aligned} \quad (11)$$

The h and v are the SM-like Higgs field and its VEV, respectively, which arise from the $SU(2)_L$ doublet $(h_1, h_2)^T$. For the charged gauge bosons, there is no mixing between W_L^\pm and W_R^\pm : $W^\pm = W_L^\pm$ and $W_H^\pm = W_R^\pm$. At $\mathcal{O}(\frac{v^2}{f^2})$, their masses and Higgs couplings are

$$\begin{aligned} m_W^2 &= \frac{1}{4}g_2^2v^2(1 - \frac{v^2}{6f^2}), & m_{W_H}^2 &= \frac{1}{2}g_2^2[\hat{f}^2 + f^2(1 - \frac{v^2}{2f^2})], \\ hWW &: \frac{1}{2}g_2^2v(1 - \frac{v^2}{3f^2}), & hW_HW_H &: -\frac{1}{2}g_2^2v(1 - \frac{v^2}{3f^2}). \end{aligned} \quad (12)$$

The neutral gauge bosons Z_H , Z and γ are linear combinations of W_L^0 , W_R^0 and W_1 . Ref. [9] gives the leading-order masses and Higgs couplings for the mass eigenstates. The diagonalization of the gauge mass matrix is performed numerically in our analysis, and the coupling of hZZ can be obtained at $\mathcal{O}(\frac{v^2}{f^2})$.

B. Mass terms of fermions

The masses of the first two generation quarks and bottom quark are obtained from the non-renormalizable operators [9]

$$\mathcal{L}_Y = \frac{y_u^{\alpha\beta}}{\Lambda}(\bar{Q}_{L\alpha}\tau_2 H_L^*)(H_R^T\tau_2 Q_{R\beta}) + \frac{y_d^{\alpha\beta}}{\Lambda}(\bar{Q}_{L\alpha}H_L)(H_R^\dagger Q_{R\beta}) + h.c., \quad (13)$$

where $\tau_2 = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}$, $Q_{L\alpha} = -i(u_{L\alpha}, d_{L\alpha})^T$ and $Q_{R\alpha} = (u_{R\alpha}, d_{R\alpha})^T$ with α being the family index. For simplicity, we assume the quark flavor mixing is small and neglect the mixing effects. From Eq. (13), we can get the Higgs boson interactions with the first two generation quarks and bottom quark:

$$\mathcal{L}_Y \simeq -\frac{y_u^\alpha}{2\Lambda}f^2s_2\bar{u}_{L\alpha}u_{R\alpha} - \frac{y_d^\alpha}{2\Lambda}f^2s_2\bar{d}_{L\alpha}d_{R\alpha} + h.c. \quad (14)$$

The mass and Higgs coupling of the quark q are given by

$$m_q = \frac{y_q}{\sqrt{2}\Lambda}fv(1 - \frac{v^2}{3f^2}), \quad h\bar{q}q : -\frac{m_q}{v}(1 - \frac{2}{3}\frac{v^2}{f^2}), \quad (15)$$

where q denotes the first two generation quarks or bottom quark.

For the lepton sector, the Yukawa interaction is similar to Eq. (13), which can generate small masses for the charged leptons and the Dirac mass terms for neutrinos.

For the top quark Yukawa interaction, in order to cancel the one-loop quadratic divergence of Higgs mass induced by the top quark, a pair of vector-like quarks (U_L, U_R) are introduced. The Lagrangian can be written as [9]

$$\mathcal{L}_t = y_L \bar{Q}_{L3} \tau_2 H_L^* U_R + y_R \bar{Q}_{R3} \tau_2 H_R^* U_L - M \bar{U}_L U_R + h.c. \quad (16)$$

where $Q_{L3} = -i(u_{L3}, d_{L3})^T$ and $Q_{R3} = (u_{R3}, d_{R3})^T$. Under left-right symmetry, $y_L = y_R = y$. From Eq.(16), we can get Higgs interaction as

$$\mathcal{L}_t \simeq -y f s_1 \bar{u}_{L3} U_R - y f c_1 \bar{u}_{R3} U_L - M \bar{U}_L U_R + h.c. \quad (17)$$

By diagonalizing the mass matrix in Eq. (17), we obtain the mass eigenstates for the top quark and heavy top quark partner T . The field t_L and T_L (t_R and T_R) are the linear combination of u_{L3} and U_L (u_{R3} and U_R), respectively. The masses and Higgs couplings of the mass eigenstates are given by [9]

$$\begin{aligned} m_t^2 &= \frac{1}{2}(M^2 + y^2 f^2 - N_t), & m_T^2 &= \frac{1}{2}(M^2 + y^2 f^2 + N_t), \\ h \bar{t} t &: -\frac{m_t}{v} C_L C_R, & h \bar{T} T &: -\frac{y}{\sqrt{2}}(S_R S_L - C_L C_R x). \end{aligned} \quad (18)$$

where

$$\begin{aligned} S_L &= \frac{1}{\sqrt{2}} \sqrt{1 - (y^2 f^2 \cos 2x + M^2)/N_t}, & C_L &= \sqrt{1 - S_L^2}, \\ S_R &= \frac{1}{\sqrt{2}} \sqrt{1 - (y^2 f^2 \cos 2x - M^2)/N_t}, & C_R &= \sqrt{1 - S_R^2}, \\ N_t &= \sqrt{(y^2 f^2 + M^2)^2 - y^4 f^4 \sin^2 2x}, \end{aligned} \quad (19)$$

with $x = \frac{v}{\sqrt{2}f}$.

C. Mass term of dark matter

In addition to the Coleman-Weinberg potential arising from gauge boson contributions, the soft left-right symmetry breaking terms, so called “ μ -term”, can give masses for \hat{h}_1^\pm and \hat{h}_2^0 [9]:

$$V_\mu = -\mu_r^2 (H_R^\dagger \hat{H}_R + h.c.) + \hat{\mu}^2 \hat{H}_L^\dagger \hat{H}_L. \quad (20)$$

In order not to reintroduce fine tuning, μ_r should be less than about $f/4\pi$. It is natural for $\hat{\mu}$ not to be much larger than f . The masses of \hat{h}_2^0 and \hat{h}_1^\pm are

$$M_{\hat{h}_2}^2 = \frac{3}{16\pi^2} \left[\frac{g_2^2}{2} (\mathcal{Z}(M_W) - \mathcal{Z}(M_{W_H})) + \frac{2g_1^2 + g_2^2}{4} \frac{M_{W_H}^2 - M_W^2}{M_{Z_H}^2 - M_Z^2} (\mathcal{Z}(M_Z) - \mathcal{Z}(M_{Z_H})) \right] + \mu_r^2 \frac{f}{\hat{f}} \cos x + \hat{\mu}^2, \\ M_{\hat{h}_1}^2 \simeq M_{\hat{h}_2}^2, \quad (21)$$

where $\mathcal{Z}(x) = -x^2(\ln \frac{\Lambda^2}{x^2} + 1)$, and the cut-off scale Λ is typically taken to be $4\pi f$. We neglect the small mass splitting between \hat{h}_2^0 and \hat{h}_1^\pm due to the electromagnetic interactions. Note that $\hat{\mu}^2$ could have either sign, which can allow us to vary the masses of \hat{h}_2^0 and \hat{h}_1^\pm as a free parameter.

The complex scalar \hat{h}_2^0 can be written as

$$\hat{h}_2^0 = \frac{\hat{S} + i\hat{A}}{\sqrt{2}}, \quad (22)$$

where \hat{S} and \hat{A} are the scalar and pseudoscalar fields, respectively. We can introduce a new quartic potential term to get the mass splitting between \hat{S} and \hat{A} , as well as their Higgs couplings [10]:

$$V_H = -\frac{\lambda_5}{2} [(H_L^\dagger \hat{H}_L)^2 + h.c.]. \quad (23)$$

From Eq. (23), we can get

$$\delta m_{\hat{S}}^2 = -\frac{\lambda_5}{2} v^2 (1 - \frac{v^2}{6f^2}), \quad \delta m_{\hat{A}}^2 = \frac{\lambda_5}{2} v^2 (1 - \frac{v^2}{6f^2}), \\ h\hat{S}\hat{S} : \lambda_5 v (1 - \frac{v^2}{3f^2}), \quad h\hat{A}\hat{A} : -\lambda_5 v (1 - \frac{v^2}{3f^2}). \quad (24)$$

Since the quartic terms $-(\lambda_5/4)h^2\hat{S}^2$ and $(\lambda_5/4)h^2\hat{A}^2$ induced by Eq. (23) have opposite sign, the one-loop quadratic divergence of Higgs mass from the \hat{S} loop and from the \hat{A} loop can be cancelled. Therefore, it is safe to take $\lambda_5 \sim 1$.

There is also a quartic term which can potentially introduce a mass splitting between \hat{h}_2^0 and \hat{h}_1^\pm [10]:

$$V'_H = \lambda_4 |\hat{H}_L^\dagger H_L|^2. \quad (25)$$

However, unlike Eq. (23), Eq. (25) can produce a dangerous contribution to the Higgs mass if λ_4 is too large, which requires $|\lambda_4| \leq \frac{1}{16\pi^2}$ with $\Lambda = 4\pi f$. Therefore, compared with Eq.

(24), the corrections of Eq. (25) to the \hat{h}_2^0 mass and Higgs coupling can be neglected. We define two parameters:

$$\begin{aligned}\delta_2 &\equiv m_{\hat{A}} - m_{\hat{S}} = \frac{\lambda_5}{(m_{\hat{A}} + m_{\hat{S}})} v^2 \left(1 - \frac{v^2}{6f^2}\right), \\ \delta_1 &\equiv m_{\hat{h}_1} - m_{\hat{S}} = \frac{\lambda_5}{2(m_{\hat{h}_1} + m_{\hat{S}})} v^2 \left(1 - \frac{v^2}{6f^2}\right).\end{aligned}\quad (26)$$

From Eqs. (21), (24) and (26), we can get the relation $\delta_2 \approx 2\delta_1$ when $m_{\hat{S}}$ is much larger than δ_2 . Actually, we checked that for a value of $m_{\hat{S}}$ not much larger than δ_2 , the relation $\delta_2 \approx 2\delta_1$ is still a good approximation. For example, for $(m_{\hat{S}}, \delta_2) = (34 \text{ GeV}, 24 \text{ GeV})$, $(70 \text{ GeV}, 40 \text{ GeV})$ and $(70 \text{ GeV}, 20 \text{ GeV})$, we found that δ_2/δ_1 is 1.77, 1.80 and 1.88, respectively. In our numerical calculations we assume $\delta_2 = 2\delta_1$ and we checked that the results are changed very little if we take δ_2/δ_1 differently as 1.77, 1.80 or 1.88.

\hat{S} is lighter than \hat{A} , and can be a candidate of dark matter. In addition to the Higgs couplings in Eq. (24), the Coleman-Weinberg potential can give the contributions to the couplings of $h\hat{S}\hat{S}$, $h\hat{A}\hat{A}$ and $h\hat{h}_1\hat{h}_1$. These expressions are complicated and can be found in [11]. In our analysis, these contributions are considered.

III. DARK MATTER SCATTERING ON NUCLEON AND HIGGS DECAY

In LRTH model, the neutral \hat{S} is a candidate for WIMP dark matter. Ref. [10] shows that there are two distinctive mass regions for \hat{S} which can give a relic density in the WMAP 3σ range: (i) low mass region, and (ii) high mass region. In this paper we focus on the low mass region where the invisible decay $h \rightarrow \hat{S}\hat{S}$ can be open, which can change other decay branching ratios and thus affect the strategy of searching for the Higgs boson at high energy colliders. For such a low mass region of \hat{S} , Fig. 9 in Ref. [10] shows that in the region satisfying the constraints of Γ_Z and WMAP 3σ relic density, δ_2 can vary in the range of 20 GeV and 40 GeV for $m_{\hat{S}} \approx 70 \text{ GeV}$, and for $30 \text{ GeV} \lesssim m_{\hat{S}} \lesssim 70 \text{ GeV}$ the value of δ_2 is around 20 GeV.

In LRTH model, the elastic scattering of \hat{S} on a nucleus receives the dominant contributions from the Higgs boson exchange diagrams. The spin-independent cross section between \hat{S} and the nucleon is given by [12]

$$\sigma_{\hat{S}p(n)}^{SI} = \frac{m_{p(n)}^2}{4\pi (m_{\hat{S}} + m_{p(n)})^2} [f^{p(n)}]^2, \quad (27)$$

where

$$f^{p(n)} = \sum_{q=u,d,s} f_{T_q}^{p(n)} \mathcal{C}_{\hat{S}q} \frac{m_{p(n)}}{m_q} + \frac{2}{27} f_{T_g}^{p(n)} \sum_{q=c,b,t} \mathcal{C}_{\hat{S}q} \frac{m_{p(n)}}{m_q}, \quad (28)$$

with [13]

$$\begin{aligned} f_{T_u}^{(p)} &\approx 0.020, & f_{T_d}^{(p)} &\approx 0.026, & f_{T_s}^{(p)} &\approx 0.118, & f_{T_g}^{(p)} &\approx 0.836, \\ f_{T_u}^{(n)} &\approx 0.014, & f_{T_d}^{(n)} &\approx 0.036, & f_{T_s}^{(n)} &\approx 0.118, & f_{T_g}^{(n)} &\approx 0.832, \\ \mathcal{C}_{\hat{S}q} &= \frac{g_{h\hat{S}\hat{S}} g_{h\bar{q}q}}{m_h^2}. \end{aligned} \quad (29)$$

Here $g_{h\hat{S}\hat{S}}$ and $g_{h\bar{q}q}$ are the couplings of $h\hat{S}\hat{S}$ and $h\bar{q}q$. Note that $\sigma_{\hat{S}p}^{SI} \approx \sigma_{\hat{S}n}^{SI}$.

In this model the major decay modes of the Higgs boson are the SM-like ones: $h \rightarrow f\bar{f}$ (SM fermion pair), WW and ZZ . The LRTH model gives corrections to these decay modes via the corresponding modified Higgs couplings

$$\Gamma(h \rightarrow XX) = \Gamma(h \rightarrow XX)_{SM} (g_{hXX} / g_{hXX}^{SM})^2, \quad (30)$$

where XX denotes fermion pairs, WW or ZZ . $\Gamma(h \rightarrow XX)_{SM}$ is the decay width in the SM, and g_{hXX} and g_{hXX}^{SM} are the couplings of hXX in the LRTH model and SM, respectively. In our calculations the relevant higher order QCD and electroweak corrections are considered using the code Hdecay [14]. In addition to the SM-like decay modes, the Higgs boson has some new important decay modes which are kinematically allowed in some parameter space: $h \rightarrow \hat{S}\hat{S}$, $h \rightarrow \hat{A}\hat{A}$ and $h \rightarrow \hat{h}_1\hat{h}_1$, whose partial widths are given by

$$\begin{aligned} \Gamma(h \rightarrow \hat{S}\hat{S}) &= \frac{g_{h\hat{S}\hat{S}}^2}{32\pi m_h} \sqrt{1 - x_{\hat{S}}}, \\ \Gamma(h \rightarrow \hat{A}\hat{A}) &= \frac{g_{h\hat{A}\hat{A}}^2}{32\pi m_h} \sqrt{1 - x_{\hat{A}}}, \\ \Gamma(h \rightarrow \hat{h}_1\hat{h}_1) &= \frac{g_{h\hat{h}_1\hat{h}_1}^2}{16\pi m_h} \sqrt{1 - x_{\hat{h}_1}}, \end{aligned} \quad (31)$$

where $x_s = 4m_s^2/m_h^2$ with $s = \hat{S}$, \hat{A} and \hat{h}_1 .

In our calculations, the free parameters involved are f , Λ , M , μ_r , $m_{\hat{S}}$ and δ_2 . \hat{f} can be determined by f , Λ , M and μ_r by requiring that the SM Higgs obtains an electroweak symmetry breaking VEV of 246 GeV [9]. The Higgs mass depends on f , Λ , M and μ_r . Note that in our numerical calculations we used the exact expressions (not only keeping the leading terms of v^2/f^2). The Z-pole precision measurements, low energy neutral current process and high energy precision measurements off the Z-pole can give strong constraints

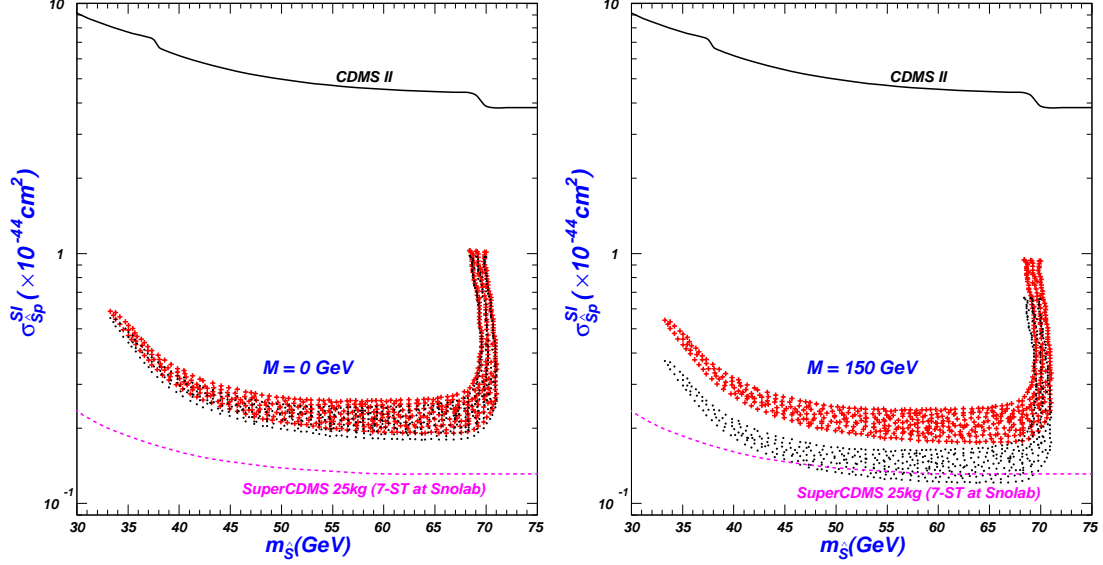


FIG. 1: Scatter plots of $\sigma_{\hat{S}p}^{SI}$ versus $m_{\hat{S}}$ from the scan over the parameters $m_{\hat{S}}$ and δ_2 in the region satisfying the constraints of Γ_Z and WMAP 3σ relic density [10]. The lower region denoted by bullets (black) and the upper region denoted by crosses (red) are for $f = 500 \text{ GeV}$ and $f = 1 \text{ TeV}$, respectively.

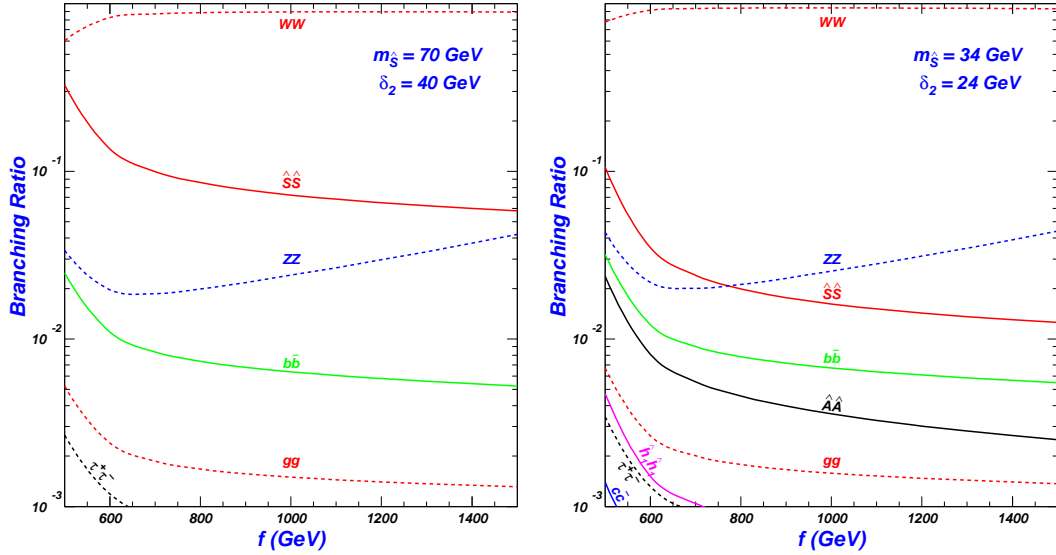


FIG. 2: The Higgs decay branching ratios versus the scale f for $M = 0 \text{ GeV}$.

on f , requiring approximately $f > 500 \text{ GeV}$. Of course, the fine tuning becomes severe if f is too large [9]. Following Ref. [11], we take the typical parameter space: $500 \text{ GeV} \leq f \leq 1500 \text{ GeV}$, $\Lambda = 4\pi f$, $\mu_r = 50 \text{ GeV}$ and $M = 0$ (150) GeV , where the Higgs mass is approximately in the range of 160 GeV and 180 GeV .

We scan over $m_{\hat{S}}$ and δ_2 in the region satisfying the constraints of WMAP 3σ relic density

and Γ_Z . The scatter plots of $\sigma_{\hat{S}p}^{SI}$ versus $m_{\hat{S}}$ are displayed in Fig. 1. We see that $\sigma_{\hat{S}p}^{SI}$ is well below the CDMS II upper bound for the low mass region of \hat{S} . The cross section can reach $1.0 \times 10^{-44} \text{ cm}^2$ for $m_{\hat{S}} \approx 70 \text{ GeV}$ where $\delta_2 = 40 \text{ GeV}$ is allowed. The coupling of $h\hat{S}\hat{S}$ increases with δ_2 , which can enhance $\sigma_{\hat{S}p}^{SI}$ sizably. The parameters M and f can have some effects on $\sigma_{\hat{S}p}^{SI}$ by changing the Higgs mass which can suppress $\sigma_{\hat{S}p}^{SI}$. Besides, the couplings of $h\hat{S}\hat{S}$ and $hq\bar{q}$ increase with f (see Eq. (15) and Eq. (24)), which can give contributions to enhance $\sigma_{\hat{S}p}^{SI}$.

In Fig. 1 we also show the projected sensitivity of SuperCDMS [15]. We see that the LRTH prediction is accessible at SuperCDMS (25kg).

In Fig. 2 we plot the Higgs decay branching ratios versus the scale f for $M = 0 \text{ GeV}$ (note that the Higgs mass can be determined by the value of f , e.g., $m_h = 159.3 \text{ GeV}$, 172.6 GeV and 178.4 GeV for $f = 500 \text{ GeV}$, 1 TeV , 1.5 TeV , respectively). The left panel shows that $Br(h \rightarrow \hat{S}\hat{S})$ is subdominant, and the largest value can reach 32% for $m_{\hat{S}} = 70 \text{ GeV}$, $\delta_2 = 40 \text{ GeV}$ and $f = 500 \text{ GeV}$. The right panel shows that the new decay modes $h \rightarrow \hat{A}\hat{A}$ and $h \rightarrow \hat{h}_1\hat{h}_1$ can be open for low values of $m_{\hat{S}}$ and δ_2 , but their decay branching ratios are relatively small. The branching ratios of all these three new decay modes decrease sizably as f increases. The reason is that the Higgs mass increases with f , and the decay width of $h \rightarrow WW$ becomes dominant for the large Higgs mass. The parameter M can have some effects on the Higgs decay modes mainly via changing the Higgs mass, which are not shown here. Besides, M can control the couplings $ht\bar{t}$ and $hT\bar{T}$, which give the dominant contributions to the decay $h \rightarrow gg$. We will show the dependence of the decay $h \rightarrow gg$ on M later.

Fig. 3 shows the scatter plots of $Br(h \rightarrow \hat{S}\hat{S})$ versus $\sigma_{\hat{S}p}^{SI}$ for $M = 0 \text{ GeV}$ and $M = 150 \text{ GeV}$, respectively. We can see that $Br(h \rightarrow \hat{S}\hat{S})$ is strongly correlated with $\sigma_{\hat{S}p}^{SI}$. When $\sigma_{\hat{S}p}^{SI}$ increases, the corresponding $Br(h \rightarrow \hat{S}\hat{S})$ also becomes large. So the SuperCDMS can probe Higgs decay $h \rightarrow \hat{S}\hat{S}$ via measuring the spin-independent WIMP-nucleon cross section, which is complementary to the exploration of Higgs boson at high energy colliders.

In Fig. 4 we plot $Br(h \rightarrow WW)$, $Br(h \rightarrow ZZ)$ and $Br(h \rightarrow b\bar{b})$ normalized to the SM predictions for several values of f . We see that the deviation from the SM prediction for each decay mode is sensitive to δ_2 , and becomes more sizable as δ_2 increases. The corrections to $Br(h \rightarrow WW)$ and $Br(h \rightarrow ZZ)$ are almost equal. For $M = 0 \text{ GeV}$, $f = 500 \text{ GeV}$, $m_{\hat{S}} = 70 \text{ GeV}$ and $\delta_2 = 40 \text{ GeV}$, the deviations for the decays $h \rightarrow VV$ ($V = W, Z$) and $h \rightarrow b\bar{b}$

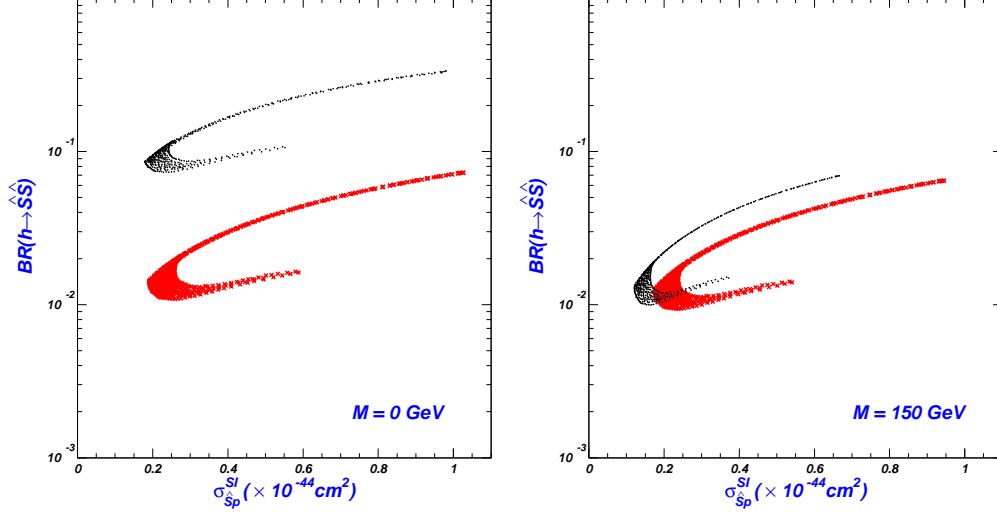


FIG. 3: Same as Fig.1, but projected on the plane of $BR(h \rightarrow \hat{S}\hat{S})$ versus σ_{Sp}^{SI} .

can be over 30% and 47.5%, respectively. The deviations from the SM predictions are also sensitive to M and f which can change Higgs mass.

In LRTH model the Higgs mass is typically in the range of 160–180 GeV and $BR(h \rightarrow \gamma\gamma)$ is severely suppressed. For $\Gamma(h \rightarrow \gamma\gamma)$, the W -boson contributions dominate over the top quark contributions [16]. The extra fermions and bosons can give some relatively small corrections, and their contributions tend to cancel each other. Therefore, the modified coupling hWW will give the dominant corrections to $\Gamma(h \rightarrow \gamma\gamma)$, and the suppression of $Br(h \rightarrow \gamma\gamma)$ approximately equals to that of $Br(h \rightarrow WW)$, which happens in littlest Higgs model with T-parity [17].

IV. HIGGS PRODUCTION AT LHC

The Higgs production at the LHC is dominated by the gluon-gluon fusion process. In the SM, the main contributions are from the top quark loop, and the LRTH model can give corrections via the modified coupling of $h\bar{t}t$ and the heavy T-quark loop. The hadronic cross section $\sigma(gg \rightarrow h)$ has a strong correlation with the decay width $\Gamma(h \rightarrow gg)$:

$$\begin{aligned} \sigma(gg \rightarrow h) &= \hat{\sigma}(gg \rightarrow h) \tau_0 \int_{\tau_0}^1 \frac{dx}{x} f_g(x, \mu_F^2) f_g\left(\frac{\tau_0}{x}, \mu_F^2\right), \\ \hat{\sigma}(gg \rightarrow h) &= \Gamma(h \rightarrow gg) \frac{\pi^2}{8m_h^3}, \end{aligned} \quad (32)$$

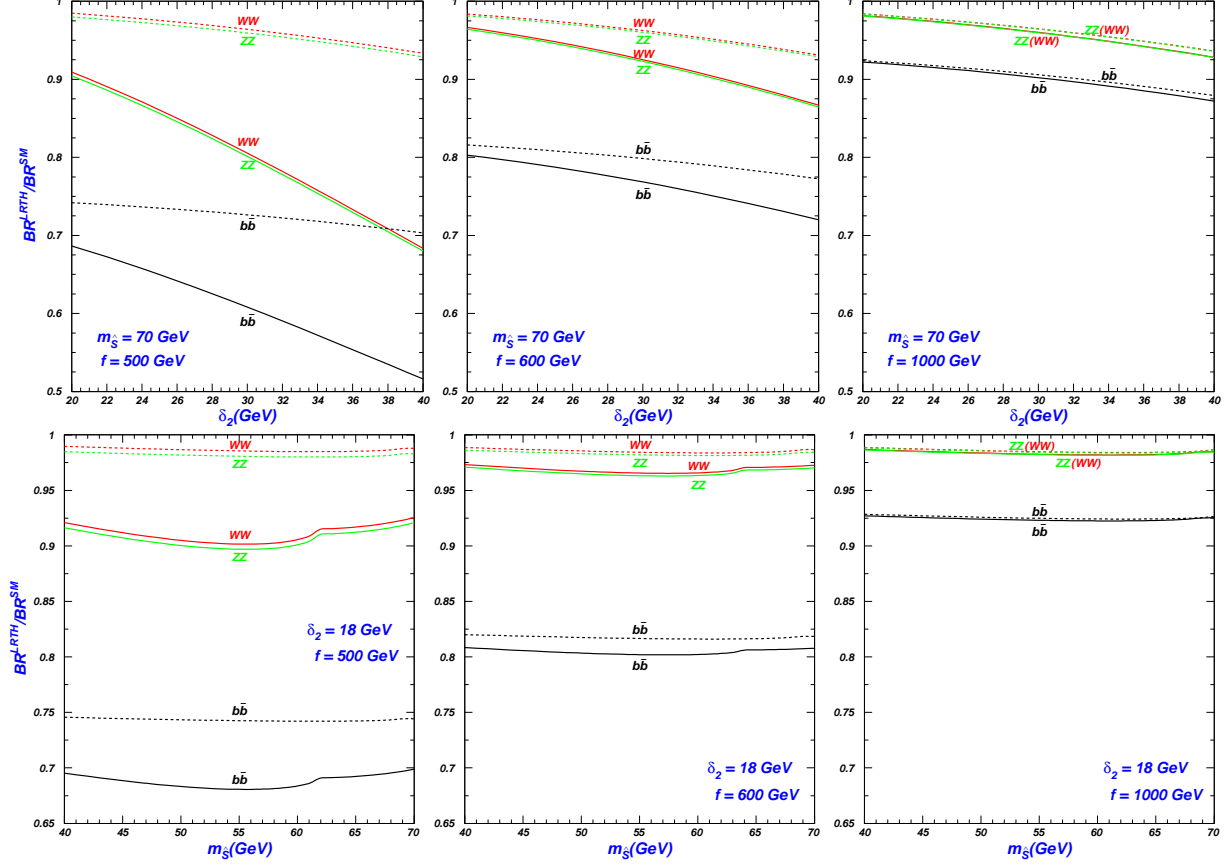


FIG. 4: The decay branching ratios $BR(h \rightarrow WW)$, $BR(h \rightarrow ZZ)$ and $BR(h \rightarrow b\bar{b})$ normalized to the SM predictions. The solid and dashed curves are for $M = 0$ GeV and $M = 150$ GeV, respectively.

where $\tau_0 = m_h^2/s$ with \sqrt{s} being the center-of-mass energy of the LHC. The parton distribution of gluon f_g is generated by CTEQ6L [18], and both the renormalization scale μ_R and the factorization scale μ_F are taken as m_h . From Eq. (32) we get

$$\frac{\sigma^{LRTH}(gg \rightarrow h)}{\sigma^{SM}(gg \rightarrow h)} = \frac{\Gamma^{LRTH}(h \rightarrow gg)}{\Gamma^{SM}(h \rightarrow gg)}. \quad (33)$$

The width $\Gamma^{LRTH}(h \rightarrow gg)$ is independent of the parameters $m_{\hat{S}}$ and δ_2 . In Fig. 5, we plot the ratio $\Gamma^{LRTH}(h \rightarrow gg)/\Gamma^{SM}(h \rightarrow gg)$ versus the scale f for $M = 0$ GeV and $M = 150$ GeV, respectively. We see that, compared with the SM prediction, the LRTH model can suppress the partial width sizably for a small value of f . As f gets large, the suppression is weakened. The deviation from the SM prediction is also sensitive to the mixing parameter M . For $f = 500$ GeV, the suppression of SM prediction can reach 27% and 32.5% for $M = 0$ GeV and $M = 150$ GeV, respectively.

Due to the suppression of the Higgs couplings with gauge bosons and top quark, the

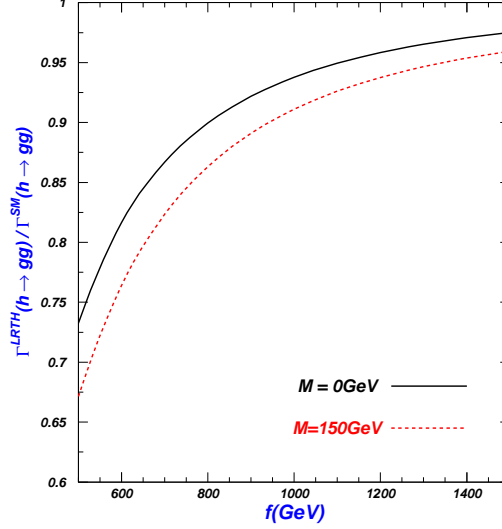


FIG. 5: The ratio $\Gamma^{LRTH}(h \rightarrow gg)/\Gamma^{SM}(h \rightarrow gg) = \sigma^{LRTH}(gg \rightarrow h)/\sigma^{SM}(gg \rightarrow h)$ versus f .

Higgs boson production rates via the weak boson fusion or in association with a pair of top quarks are suppressed. Although $Br(h \rightarrow WW)$ can be sizably suppressed in some parameter space, the decay mode $h \rightarrow WW$ is still an excellent channel for searching for Higgs boson with an intermediate mass [19].

In Table 1 the ratio $\frac{\sigma^{LRTH} \times BR^{LRTH}}{\sigma^{SM} \times BR^{SM}}$ is listed for each main channel: $gg \rightarrow h$, $VV \rightarrow h$ or $pp \rightarrow h\bar{t}t$ followed by $h \rightarrow VV$ or $h \rightarrow b\bar{b}$. Table 1 shows that such a ratio of events can be sizably suppressed in the LRTH model and the suppression effects could exceed the experimental uncertainty (10%-20%) [20]. Therefore, it is possible to probe the LRTH model via these Higgs production modes at the LHC.

Note that similar exotic decays for the SM-like Higgs boson may also be predicted by some other new physics models like the little Higgs models and supersymmetric or two Higgs-doublet models [21]. A common feature of their phenomenology is the suppression of the conventional visible channels of the Higgs boson. To distinguish between different models, all the channels of Higgs production should be jointly analyzed and a linear collider is an ideal machine for such a purpose [22].

Recently both CDF and D0 Collaborations at the Tevatron have searched for the SM Higgs boson and excluded the mass range between 162 GeV and 166 GeV at 95% CL [23], using all significant production modes: $gg \rightarrow h$, $VV \rightarrow h$, $q\bar{q} \rightarrow Vh$ ($V = W, Z$) followed by $h \rightarrow WW$. The productions $gg \rightarrow h$ and $VV \rightarrow h$ are suppressed, similar to the case at the LHC shown in Table 1. The productions $q\bar{q} \rightarrow Vh$ are also suppressed, just like $VV \rightarrow h$,

TABLE I: The ratio $\frac{\sigma^{LRTH} \times BR^{LRTH}}{\sigma^{SM} \times BR^{SM}}$ for various channels ($gg \rightarrow h$, $VV \rightarrow h$ or $pp \rightarrow h\bar{t}t$ followed by $h \rightarrow VV$ or $h \rightarrow b\bar{b}$) with $m_{\hat{S}} = 70$ GeV and $\delta_2 = 40$ GeV.

	$h \rightarrow b\bar{b}$			$h \rightarrow VV$		
	$f=500$ GeV	$f=600$ GeV	$f=1$ TeV	$f=500$ GeV	$f=600$ GeV	$f=1$ TeV
$gg \rightarrow h$ ($M = 0$ GeV)	0.38	0.59	0.82	0.50	0.71	0.87
($M = 150$ GeV)	0.47	0.59	0.80	0.63	0.71	0.85
$VV \rightarrow h$ ($M = 0$ GeV)	0.46	0.66	0.85	0.60	0.80	0.90
($M = 150$ GeV)	0.62	0.71	0.85	0.82	0.85	0.91
$pp \rightarrow t\bar{t}h$ ($M = 0$ GeV)	0.50	0.70	0.87	0.66	0.85	0.92
($M = 150$ GeV)	0.59	0.69	0.85	0.79	0.83	0.90

because these productions are induced by the suppressed couplings hWW and hZZ . Since all these productions $gg \rightarrow h$, $VV \rightarrow h$ and $q\bar{q} \rightarrow Vh$ followed by $h \rightarrow WW$ can be severely suppressed in the LRTH model, the constraints on the Higgs mass from the Tevatron will be largely relaxed in the LRTH model.

V. CONCLUSION

In LRTH model, the scalar \hat{S} is a natural candidate for WIMP dark matter, and the Higgs boson mass is typically in the range of 160 - 180 GeV. Since the invisible decay $h \rightarrow \hat{S}\hat{S}$ can affect other decay branching ratios, and also has a strong correlation with the scattering on nucleon, we in this work focused on the low mass region of \hat{S} so that the decay $h \rightarrow \hat{S}\hat{S}$ can be open. We obtained the following observations: (i) The cross section of \hat{S} scattering on nucleon can naturally satisfy the CDMS II upper bound, and can be large enough to be accessible at SuperCDMS; (ii) The Higgs boson can have a sizable invisible decay $h \rightarrow \hat{S}\hat{S}$, whose branching ratio can reach 32% and has a strong correlation with the cross section of \hat{S} scattering on nucleon. However, the branching ratios of other new decay modes $h \rightarrow \hat{A}\hat{A}$ and $h \rightarrow \hat{h}_1\hat{h}_1$ are small; (iii) The branching ratios of the conventional decay modes of the Higgs boson, $h \rightarrow VV$ ($V = W, Z$) and $h \rightarrow b\bar{b}$, can be suppressed over 30% and 47.5%, respectively; (iV) The Higgs production cross sections times the branching ratios of the

conventional decays can be all sizably suppressed. So, it is possible to probe the LRTH model via the Higgs productions at the LHC.

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